Confinement and short distance physics

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Abstract

We consider non-perturbative effects at short distances in theories with confinement. The analysis is straightforward within the Abelian models in which the confinement arises on classical level. In all cases considered (compact $U(1)$ in $3D$ and $4D$, dual Abelian Higgs model) there are non–perturbative contributions associated with short distances which are due to topological defects. In QCD case, both classical and quantum effects determine the role of the topological defects and the theoretical analysis has not been completed so far. Generically, the topological defects would result in $1/Q^2$ corrections going beyond the standard Operator Product Expansion. We review existing data on the power corrections and find that the data favor existence of the novel corrections, at least at the mass scale of (1-2) GeV. We indicate crucial experiments which could further clarify the situation on the phenomenological side.

1. The perturbative QCD describes basic features of hard processes, i.e. processes characterized by a large mass scale Q. On the other hand, the perturbative QCD does not encode the effects of confinement. Hence there is a growing interest in power corrections to the parton model which might be sensitive to the nature of confinement (for a recent review see [1]). Moreover, in the the Abelian Higgs model the leading power correction at short distances does reflect confinement [2]. Namely, there exist short strings which are responsible for a stringy potential between confined charges at vanishing distances $r \to 0$.

In this note we are exploring Abelian models with confinement in a more regular way by including into consideration the compact three and four dimensional $U(1)$ theory. In all the cases the confinement mechanism can be understood classically. Moreover, we find that the confinement results in additional terms in the static potential at short distances. How topological defects can be manifested at short distances, is easy to understand on the example of the Dirac string. Naively, its energy diverges quadratically in the ultraviolet but in compact $U(1)$ it is normalized to zero [3] changing the power corrections at short distances. As for non-Abelian theories, the Dirac strings are also allowed because of the compactness of the corresponding $U(1)$ subgroup. In this sense, there is a similarity between non–Abelian and Abelian cases. However, there is an important difference as well. On the classical level, the Dirac strings may end up with monopoles which have a vanishing non-Abelian action. Thus,

the monopoles observed within the $U(1)$ projection of QCD (for a review see, e.g., [4]) is a result of an interplay between classical and quantum effects. In all the generality, one may say that the topological defects in QCD are marked rather by singular potentials than by a large non-Abelian action. Singular gauge potentials might be artifact of the gauge fixing and it is not a priori clear whether they can result in physical effects. Therefore, we will turn at this point to analysis of existing data on the power corrections. The data seem clearly favor the novel $1/Q^2$ corrections.

2. First, we will outline very briefly the standard approach to the power corrections which allows to account for soft non-perturbative field configurations (for further references see, e.g., [1]). Consider first a QED example [5]. Namely, let an e^+e^- pair be placed at distance r near the center of a conducting cage of size L, $L \gg r$. Then the potential energy of the pair can be approximated as

$$V_{e\bar{e}}(r) \approx -\frac{\alpha_e}{r} \; + \; const \cdot \frac{\alpha_e r^2}{L^3}, \qquad L \gg r \tag{1}$$

and the second term is a power correction to the Coulomb interaction. The derivation of (1) is of course straightforward classically, since the correction is nothing else but interaction of the dipole with its images. In the QCD case, one concludes by analogy that the heavy quark potential at short distances looks as (for explanations and further references see, e.g. [6]):

$$\lim_{r \to 0} V_{Q\bar{Q}}(r) \; = \; -\frac{c_{-1}}{r} \; + \; const \cdot \Lambda_{QCD}^3 \, r^2, \tag{2}$$

where c_{-1} is calculable perturbatively as a series in α_s. Note the absence of a linear correction to the potential at short distances.

On the other hand, Eq. (1) can be derived also in terms of one-photon exchange. The power correction is related then to a change in modes of the electromagnetic field confined in the cage as compared to the case of the infinite space. The change is of order unit at frequencies $\omega \sim 1/L$. Similarly, the logic behind Eq. (2) is that the perturbative gluon propagator is modified strongly by at $\omega \sim \Lambda_{QCD}^{-1}$. In case of other processes, the relevant Feynman graphs can be more complicated of course. The power corrections still correspond to the infrared sensitive part of Feynman propagators which are obviously modified by the physics of large distances. The Operator Product Expansion (OPE) allows for a regular way to parameterize such corrections (for a review see [1]).

3. Intuitively, the power corrections in QCD could be very different from the conducting cage case discussed above. Indeed color particles produced at short distances find no cage but rather build up the confining field configuration in the course of the interaction between themselves and with the vacuum. The complicated space-time picture of interaction in confining theories was studied by Gribov [7]. Thus, it could be instructive to analyze the effects of the confinement at short distances in some simple models.

The first example of a theory where the OPE does not work in fact goes back to the paper in Ref. [3]. However, since it has not been discussed in connection with the OPE, we will explain this example in some detail. The action is that of free photons:

$$S \; = \; \frac{1}{4e^2} \int d^4x F_{\mu\nu}^2 \tag{3}$$

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where $F_{\mu\nu}$ is the field strength tensor, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. Although the theory looks absolutely trivial, it is not the case if one admits Dirac strings into the theory. Naively, the energy associated with the Dirac strings is infinite:

$$E_{\text{Dirac string}} = \frac{1}{8\pi} \int d^3r \, \mathbf{H}^2 \sim l \cdot A \left(\frac{\text{magnetic flux}}{A} \right)^2 \to \infty \qquad (4)$$

where l, A are the length and area of the string, respectively. Since the magnetic flux carried by the string is quantized and finite the energy diverges quadratically in the ultraviolet, i.e. in the limit $A \to 0$. However within the lattice regularization the action of the string is in fact zero because of the compactness of the $U(1)$, Ref. [3].

Now, the Dirac strings may end up with monopoles. The action associated with the monopoles diverges in ultraviolet,

$$\int \frac{d^3r}{8\pi} \mathbf{H}^2 \sim \frac{1}{e^2 a} \qquad (5)$$

where a is a (small) spatial cut off. If the length of a closed monopole trajectory is L, then the suppression of such a configuration due to a non-vanishing action is of order

$$e^{-S} \sim \exp\left(-const \, L/e^2 \right). \qquad (6)$$

On the other hand, there are different ways to organize a loop of length L. This is the entropy factor. It is known to grow exponentially with L as $\sim \exp(\, const' L \,)$. At some $e^2_{crit} \sim 1$ there is a phase transition to the monopole condensation.

The potential between external test charges is Coulombic at all the distances for $e^2 < e^2_{crit}$ and linear for $e^2 > e^2_{crit}$. Since there are no perturbative graphs at all in the theory with the action (3) this phenomenon clearly goes beyond the OPE. In this case, however, the violation of the OPE is too strong. Indeed, the lattice spacing a is the only dimensional parameter of the problem. As a result Coulomb potential is not simply modified by linear corrections but rather eliminated for $e^2 > e^2_{crit}$ at all the distances.

4. Consider next the Dual Abelian Higgs Model with the action

$$S = \int d^4x \left\{ \frac{1}{4g^2} F_{\mu\nu}^2 + \frac{1}{2}|(\partial - iA)\Phi|^2 + \frac{1}{4}\lambda(|\Phi|^2 - \eta^2)^2 \right\}, \qquad (7)$$

here g is the magnetic charge, $F_{\mu\nu} \equiv \partial_\mu A_\nu - \partial_\nu A_\mu$. The gauge boson and the Higgs are massive, $m_V^2 = g^2\eta^2$, $m_H^2 = 2\lambda\eta^2$. There is a well known Abrikosov-Nielsen-Olesen (ANO) solution to the corresponding equations of motion. The dual ANO string may end up with electric charges. As a result, the potential for a test charge–anti-charge pair grows linearly at large distances:

$$V(r) = \sigma_\infty r , \qquad r \to \infty. \qquad (8)$$

Note that there is a Dirac string resting along the axis of the ANO string connecting monopoles and its energy is still normalized to zero.

An amusing effect occurs if one goes to distances much smaller than the characteristic mass scales $m_{V,H}^{-1}$. Then the ANO string is peeled off and one deals with a naked (dual) Dirac string. The manifestation of the string is that the Higgs field has to vanish along a line connecting the external charges. Otherwise, the energy of the Dirac string would jump to infinity anew.

As a result of the boundary condition that Φ vanishes on a line connecting the charges the potential contains a stringy piece at short distances [2]:

$$\lim_{r \to 0} V(r) = \frac{c_{-1}}{r} + \sigma_0 \cdot r . \tag{9}$$

The string tension σ_0 smoothly depends on the ratio m_H/m_V. In the Bogomol'ny limit ($m_H = m_V$) which is favored by the fits of the lattice simulations [8] the string tension

$$\sigma_0 \approx \sigma_\infty, \tag{10}$$

i.e. the effective string tension numerically is the same at all distances.

5. Consider now $3D$ compact electrodynamics. As is well known [9], the charge–anti-charge potential is then linear at large r. Below we consider the string tension σ_0 at small distances and show that it has a non-analytical piece.

As usual, it is convenient to perform the duality transformation, and work with the corresponding Sine-Gordon theory. The expectation value of the Wilson loop in dual variables is:

$$W = \frac{1}{\mathcal{Z}} \int \mathcal{D}\chi e^{-S(\chi, \eta_c)}, \tag{11}$$

where

$$S(\chi, \eta_c) = \left(\frac{e}{2\pi}\right)^2 \int d^3x \left\{ \frac{1}{2}(\vec{\partial}\chi)^2 + m_D^2(1 - \cos[\chi - \eta_c]) \right\}, \tag{12}$$

m_D is the Debye mass and $S(\chi, 0)$ is the action of the model. If static charge and anti-charge are placed at the points $(-R/2, 0)$ and $(R/2, 0)$ in the x_1, x_2 plane (x_3 is the time axis), then

$$\eta_c = \text{arctg}[\frac{x_2}{x_1 - R/2}] - \text{arctg}[\frac{x_2}{x_1 + R/2}], \qquad -\pi \le \eta_c \le \pi. \tag{13}$$

Below we present the results of the numerical calculations of the dimensionless string tension,

$$\sigma = \partial E / \partial(m_D R) , \tag{14}$$

$$E = \int d^2x \left\{ \frac{1}{2}(\vec{\partial}\chi)^2 + m_D^2(1 - \cos[\chi - \eta_c]) \right\} . \tag{15}$$

Note that the energy E is measured in the units of the dimensional factor $(\frac{e}{2\pi})^2$ (*cf.* (12)). Variation of functional (15) leads to the equation of motion $\Delta\chi = m_D^2 \sin[\chi - \eta_c]$. For finite R we can solve this equation numerically. The energy E versus $m_D R$ is shown on Fig.1(a). At large separations between the charges ($m_D R \gg 1$) it tends to the asymptotic linear behavior $E = 8 m_D R$ which can be obtained also analytically [9].

At small distances there is a contributions of Yukawa-type to the energy (15), which should be extracted explicitly. Note that in course of rewriting original $3D$ compact electrodynamics in
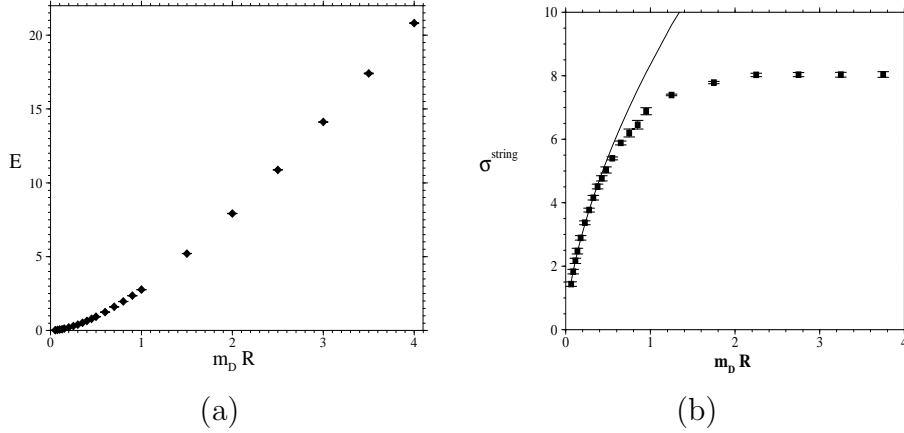
Figure 1: (a) The dimensionless string tension (14) of the charge – anti-charge separated by the distance R in 3D compact U(1) theory; (b) The string tension σ^{string} (17) as a function of $m_D R$ with corresponding best fitting function (see text).

the form (11-12) the Coulomb potential was already subtracted, so that (15) contains Yukawa-like piece without singularity at $R = 0$. It is not difficult to find the corresponding coefficient:

$$E = E^{string} - 2\pi(K_0[m_D R] + \ln[m_D R]) \tag{16}$$

where $K_0(x)$ is the modified Bessel function and E^{string} is the energy of the charge–anti-charge pair which is only due to the string formation. The corresponding string tension

$$\sigma^{string} = \sigma + 2\pi(-K_1[m_D R] + \frac{1}{m_D R}) \tag{17}$$

is shown on Fig.1(b). We found that the best fit of numerical data for small values of $m_D R$ is by the function $\sigma^{string} = const \cdot (m_D R)^\nu$ which gives $\nu \approx 0.6$.

Thus the non-analytical potential associated with small distances is softer than in the case of the Abelian Higgs model. The source of the non-analyticity is the behavior of the function $\eta_{\mathcal{C}}(x_1, x_2)$ eq. (13) which is singular along the line connecting the charges, see Fig.2(a).

6. The compact electrodynamics is usually considered as the limit of Georgi–Glashow model, when the radius of the 't Hooft – Polyakov monopole tends to zero. For a non-vanishing monopole size the problem of evaluating the potential at small distances becomes rather involved. To avoid unnecessary complications we consider the 3D Georgi–Glashow model in the BPS limit. The 't Hooft – Polyakov monopole corresponds then to the fields:

$$\Phi^a = \frac{x^a}{r}\left(\frac{1}{\tanh(\mu r)} - \frac{1}{\mu r}\right), \tag{18}$$

$$A_i^a = -\varepsilon^{aic}\frac{x^a}{r}\left(\frac{1}{r} - \frac{\mu}{\sinh(\mu r)}\right), \quad A_0^a = 0. \tag{19}$$
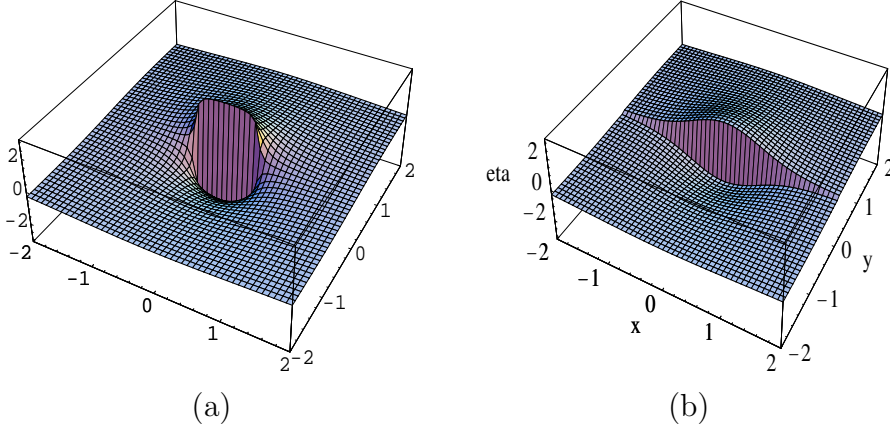
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Figure 2: Function $\eta_{\mathcal{C}}$ in compact $U(1)$ theory (a) and for extended monopoles (b), eq. (22)

The contribution of this monopole to the *full non-Abelian* Wilson loop W can be calculated analytically. If the static charges are placed at points $\pm \vec{R}/2$ in the (x_1, x_2) plane the result is:

$$W(\vec{b}_1, \vec{b}_2, \mu) = \cos h(\mu b_1) \, \cos h(\mu b_2) + \frac{(\vec{b}_1 \cdot \vec{b}_2)}{b_1 \, b_2} \sin h(\mu b_1) \, \sin h(\mu b_2) \,, \tag{20}$$

here $\vec{b}_{1,2} = \vec{x}_0 \pm \vec{R}/2$, $b_k = |\vec{b}_k|$, $\vec{x}_0$ is the center of the 't Hooft – Polyakov monopole and

$$h(x) = \frac{\pi}{2} - \frac{x}{2} \int\limits_{-\infty}^{+\infty} \frac{\mathrm{d}\zeta}{\sqrt{x^2 + \zeta^2} \, \sinh \sqrt{x^2 + \zeta^2}} \,. \tag{21}$$

One way to represent (20) in terms of the function $\eta_{\mathcal{C}}$ introduced earlier is:

$$\eta_{\mathcal{C}}(x_0, R, \mu) = \mathrm{sign}(y) \, \arccos W(\vec{b}_1, \vec{b}_2, \mu) \,. \tag{22}$$

In the limit $\mu R \to \infty$ $W(\vec{b}_1, \vec{b}_2, \mu) \to \cos \eta_{\mathcal{C}}$ and $\eta_{\mathcal{C}}(x_0, R, \mu)$ coincides with the definition (13). For small μR the function $\eta_{\mathcal{C}}$ eq. (22) is singular not only between external charges, but also outside this region (see Fig.2(b)) although the strength of singularity gets smaller. In the limit of vanishing $\eta_{\mathcal{C}}$ the string tension at small distances σ_0 apparently goes to zero.

To summarize, it is natural to expect that in the Georgi-Glashow model the non-analytical piece in the potential disappears at distances much smaller than the monopole size. Note, however, that to evaluate the potential consistently in this case one should have taken into account also the modification of the interaction due to the finite size of the monopoles.

7. Knowing the physics of the Abelian models outlined above it is easy to argue that the perturbative vacuum of QCD is not stable as well. Indeed, let us make the lattice coarser a la Wilson until the effective coupling of QCD would reach the value where the phase transition in the compact $U(1)$ occurs. Then the QCD perturbative vacuum is unstable against the monopole condensation. The actual vacuum state can of course be different from the $U(1)$ case but it

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cannot remain perturbative. Similar remark with respect to formation of Z_2 vortices was in fact made long time ago [10].

The existence of the infinitely thin topological defects in QCD makes it close akin of the Abelian models considered above. However, the non-Abelian nature of the interaction brings in an important difference as well. Namely, the topological defects in QCD are marked rather by singular potentials than by a large non-Abelian action. Consider first the Dirac string.

Introduce to this end a potential which is a pure gauge:

$$A_\mu = \Omega^{-1} \partial_\mu \Omega \tag{23}$$

and choose the matrix Ω in the form:

$$\Omega(x) = \begin{pmatrix} \cos\frac{\theta}{2} & \sin\frac{\theta}{2} \, e^{-i\varphi} \\ -\sin\frac{\theta}{2} \, e^{i\varphi} & \cos\frac{\theta}{2} \end{pmatrix} \tag{24}$$

where φ and θ are azimuthal and polar angles, respectively. Then it is straightforward to check that we generated a Dirac string directed along the x_3-axis ending at $x_3 = 0$ and carrying the color index $a = 3$. It is quite obvious that such Abelian-like strings are allowed by the lattice regularization of the theory.

The crucial point, however, is that the non-Abelian action associated with the potential (23) is identical zero. On the other hand, in its Abelian components the potential looks as a Dirac monopole, which are known to play important role in the Abelian projection of QCD (for a review see, e.g., [4]). Thus, there is a kind of mismatch between short- and large-distance pictures. Namely, if one considers the lattice size $a \to 0$, then the corresponding coupling $g(a) \to 0$ and the solution with a zero action (23) is strongly favored at short distances. At larger distances we are aware of the dominance of the Abelian monopoles which have a non-zero action. The end-points of a Dirac string still mark centers of the Abelian monopole. Thus, monopoles can be defined as point-like objects topologically in terms of singular potentials, not action.

Similar remarks hold in case of the Z_2 vortices. Namely the Z_2 vortices which have a typical size of order Λ_{QCD}^{-1} can be defined topologically in terms of the so called P-vortices which are infinitely thin but gauge dependent, see [11] and references therein. To detect the P-vortices one uses the gauge maximizing the sum

$$\sum_l |Tr \, U_l|^2 \tag{25}$$

where l runs over all the links on the lattice. The center projection is obtained by replacing

$$U_l \to \text{sign} \, (Tr \, U_l). \tag{26}$$

Each plaquette is marked either as $(+1)$ or (-1) depending on the product of the signs assigned to the corresponding links. The P-vortex then pierces a plaquette with (-1). Moreover, the fraction p of the total number of plaquettes pierced by the P-vortices and of the total number of all the plaquettes N_T, was found to obey numerically the scaling law

$$p = \frac{N_{vor}}{N_T} \sim f(\beta) \tag{27}$$

where the function $f(\beta)$ is such that p scales like the string tension. Assuming independence of the piercing for each plaquette one has then for the center-projected Wilson loop W_{cp}:

$$W_{cp} = [(1-p)(+1) + p(-1)]^A \approx e^{-2pA} \tag{28}$$

where A is the number of plaquettes in the area stretched on the Wilson loop. Numerically, Eq. (28) reproduces the full string tension.

It is quite obvious that the P-vortices, since they are constructed on negative links, correspond in the continuum limit to singular gauge potentials of order a^{-1}. Moreover, the large potentials should mostly cancel if the corresponding field-strength tensors are calculated because of the asymptotic freedom. The argumentation is essentially the same as outlined above for the monopoles, see, e.g. [12] and references therein.

At the moment, it is difficult to say a priori whether the topological defects defined in terms of singular potentials can be considered as infinitely thin from the physical point of view. They might be gauge artifacts. Phenomenologically, using the topologically defined point-like monopoles or infinitely thin P-vortices one can measure non-perturbative $Q\bar{Q}$ potential at all the distances. It is remarkable therefore that the potentials generated both by monopoles [13] and P-vortices [11] turn to be linear at all the distances measured:

$$V_{non-pert}(r) \approx \sigma_\infty r \quad \text{at all } r \tag{29}$$

Note that the Coulomb-like part is totally subtracted out through the use of the topological defects. Moreover, the no-change in the slope (29) agrees well with the dual Abelian Higgs model as discussed above (for alternative approaches see [14, 15]).

The numerical observation (29) is by no means trivial. If it were so that only the non-Abelian action counts, then the non-perturbative fluctuations labeled by the Dirac strings or by P-vortices are bulky (see discussion above) and the corresponding $Q\bar{Q}$ potentials (29) should have been quadratic at small r. This happens, for example, in the model [16] with finite thickness of Z_2 vortices. Similarly, if the lessons from the Georgi–Glashow model considered above apply the finite size of the monopoles would spoil linearity of the potential at short distances.

To summarize, direct measurements of the non-perturbative $Q\bar{Q}$ potential indicate the presence of a stringy potential at short distances. The measurements go down to distances of order $(2\text{ GeV})^{-1}$.

8. In view of the results (29) it is interesting to reexamine the power corrections with the question in mind, whether there is room for novel stringy corrections. From the dimensional considerations alone it is clear that the new corrections are of order σ_0/Q^2 where Q is a large generic mass parameter characteristic for problem in hand. Also, the ultraviolet renormalons in 4D indicate the same kind of correction, see [6] and references therein. Note that unlike the case of the non-perturbative potential discussed above, other determinations of the power corrections ask for a subtraction of the dominating perturbative part and this might make the results less definitive.

(i) The first claim of observation of the non-standard $1/Q^2$ corrections was made in Ref. [17]. Namely, it was found that the expectation value of the plaquette minus perturbation theory

contribution shows $1/Q^2$ behavior. On the other hand, the standard OPE results in a $1/Q^4$ correction.

(ii) The lattice simulation [18] do not show any change in the slope of the full $Q\bar{Q}$ potential as the distances are changed from the largest to the smallest ones where the Coulombic part becomes dominant. An explicit subtraction of the perturbative corrections at small distances from $Q\bar{Q}$ potential in lattice gluodynamics was performed in ref.[19]. This procedure gives $\sigma_0 \approx 5\sigma_\infty$ at very small distances.

(iii) There exist lattice measurements [20] of the fine splitting of $Q\bar{Q}$ levels as function of the heavy quark mass. The Voloshin-Leutwyler [21] picture predicts a particular pattern of the heavy mass dependence of this splitting. Moreover, these predictions are very different from the predictions based, say, on the Buchmuller-Tye potential [22] which adding a linear part to the Coulomb potential. The numerical results favor the linear correction to the potential at short distances.

(iv) Analytical studies of the Bethe-Salpeter equation and comparison of the results with the charmonium spectrum data favor a non-vanishing linear correction to the potential at short distances [23].

(v) The lattice-measured instanton density as a function of the instanton size ρ does not satisfy the standard OPE predictions that the leading correction is of order ρ^4. Instead, the leading corrections is in fact quadratic [25].

(vi) One of the most interesting manifestations of short strings might be the $1/Q^2$ corrections to current correlation functions $\Pi_j(Q^2)$. It is not possible to calculate the coefficient of front of the $1/Q^2$ terms from first principles, however, in Ref. [24] it was suggested to simulate this correction by a tachyonic gluon mass. On one hand, the tachyonic mass imitates the stringy piece in the potential at short distances. On the other hand, it can be used in one-loop calculations of the correlation functions. Rather unexpectedly, the use of the tachyonic gluon mass ($m_g^2 = -0.5$ GeV^2) explains well the behavior of $\Pi_j(Q^2)$ in various channels. To check the model further, it would be very important to perform accurate calculations of various correlators $\Pi_j(Q^2)$ on the lattice.

9. As seen from the points *(i)*-*(vi)* above, the existence of the novel quadratic corrections is strongly supported by the data. There are, however, two caveats to the statement that the novel short-distance power corrections have been detected. On the theoretical side, the existence of short strings has been proven only within the Abelian Higgs model. As for the QCD itself, the analysis is so far inconclusive. On the experimental side, the data always refer to a limited range of distances. In particular, the linear non-perturbative potential has been observed at distances of order of one lattice spacing which in physical units is about $(1 \div 2$ GeV$)^{-1}$. One cannot rule out that at shorter distances the behavior of the non-perturbative power corrections changes (see, e.g., [15, 25]). Which would be a remarkable phenomenon by itself.

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